Charmonium Production in Nucleus-Nucleus Collisions

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Abstract

This is a review of theoretical attempts to describe production of heavy quark bound states in nucleus-nucleus collisions, in particular, the relative suppression of J/ψ and ψ' production observed in these reactions. The review begins with a survey of experimental data for proton-induced reactions and their theoretical interpretation. The evidence for additional suppression in nucleus-nucleus collisions is discussed and various theoretical models of charmonium absorption by comoving matter are presented and analyzed. The review concludes with suggestions for future research that would help clarify the implications of J/ψ suppression in Pb + Pb collisions observed by the NA50 experiment.

1 Introduction

Ever since their discovery two decades ago, bound states of heavy quarks have served as one of the primary probes of quantum chromodynamics. The spectroscopy of charmonium and bottomonium states provided crucial tests for the theory of forces among quarks, culminating in recent years in the precise determination of the strong coupling constant α_s by means of lattice simulations.

Because of their relatively small size, bound states of heavy quarks interact only rather weakly with other hadrons. As Matsui and Satz pointed out more than a decade ago, this property makes them excellent candidates as probes of superdense hadronic matter, created in collisions between nuclei at high energy [1]. As long as this matter were composed of hadrons, so they argued, it would only weakly affect simultaneously created heavy quark bound states (J/ψ, ψ', Υ, etc.). On the other hand, if a color deconfined state, i.e. a quark-gluon plasma, were formed in the nuclear reaction the formation of heavy quark bound states would be severely suppressed, because the attractive color force between the quarks is screened by the plasma. Quantitative calculations of this effect at finite temperature

can be performed in the framework of lattice gauge theory. These calculations yield a state-specific critical temperature $T_{\mathrm{d}}^{(i)}$ above which the bound state (i) dissociates into the plasma. Whereas the charmonium states dissociate soon above the critical temperature T_{c} of the QCD phase transition $(T_{\mathrm{d}}^{(J/\psi)} \approx 1.2\, T_{\mathrm{c}})$, the bottomonium states survive up to higher temperatures $(T_{\mathrm{d}}^{(v)} \approx 2 T_{\mathrm{c}})$.

A systematic account of how these characteristic properties of heavy quark bound states may be put to use in the study of superdense matter, especially for the purpose of establishing the existence of a new deconfined phase of QCD, was presented several years ago by Karsch and Satz [2]. Since then, new theoretical insight into the elementary process of heavy quarkonium production has led to a reassessment of the suppression of charmonium production in hadronic interactions with nuclei. In fact, as will be discussed in the next section, all experimental data taken with nuclear projectiles (^{16}O, ^{32}S) up to 1994 can probably be explained as final state interactions of the produced heavy quark pair with the nucleons contained in the colliding nuclei. The discovery of enhanced suppression of J/ψ production in collisions between two Pb nuclei at the CERN-SPS in 1996, therefore, generated intense interest. Does the "anomalous" J/ψ suppression observed in Pb + Pb collisions indicate the formation of a quark-gluon plasma in these reactions? Has this new state of matter been finally detected?

In order to answer this question, all other conceivable mechanisms that could explain the observed suppression effect need to be ruled out. This is not an easy task because reactions among heavy nuclei at high energy are processes of great complexity, and there is little hope that they can be understood in all details. It is therefore necessary to allow for a considerable range of parameters in different reaction models and to exclude competing suppression mechanisms under a variety of assumed scenarios.

It is the purpose of this review to assess the status of the theoretical tools that are needed to establish the possible origin of the observed "anomalous" J/ψ suppression. The next section summarizes the current experimental information on J/ψ and ψ' suppression in proton-nucleus and nucleus-nucleus collisions. The explanation of all data (except those from Pb + Pb collisions) in terms of nuclear final state interactions will also be discussed in this section. In Section 3, a number of possible mechanisms for the enhanced suppression seen in Pb + Pb collision will be surveyed and their current theoretical uncertainties will be analyzed. Section 4 deals with the quantitative analysis of the Pb + Pb data in the framework of specific reaction models. The review closes with a list of suggestions for future research, mostly theoretical, that could help sharpen the tool of charmonium spectroscopy

for the exploration of the properties of superdense hadronic matter.

The discussion here is mostly based on publications that were available in June 1997. Later developments are briefly discussed in a separate section at the end of the review. For a survey of the state of this research field from a somewhat different perspective, the reader is urged to also consult the recent review by Kharzeev [3].

2 Phenomenology

2.1 Proton-Induced Reactions

The production of charmonium states in proton interactions with nuclei has been studied extensively at Fermilab as well as at CERN. Since these experiments observe the $\mu^+\mu^-$ decay mode of the J/ψ and ψ', they also provide information on the nuclear target dependence of the $\mu^+\mu^-$ continuum spectrum produced by light quark-antiquark annihilation (Drell-Yan process, DY). Whereas the Drell-Yan cross section is observed to grow almost exactly in proportion to the nuclear mass number A, the J/ψ and ψ' production cross sections grow less rapidly. The target dependence is often parametrized as a power law, $\sigma = \sigma_0 A^\alpha$. A value $\alpha <$ 1 indicates "suppression" of charmonium production, compared with the naive expectation $\sigma/A = $ const. obtained when one neglects final state interactions. The argument is that the strictly linear A-dependence of the DY cross section rules out initial state interactions as the origin of the reduced production of charmonium states.[1]

The A-dependence of the nuclear J/ψ and ψ' cross sections is remarkably similar. Both can be fit by the same exponent $\alpha \approx 0.92$, as shown in Fig. 1. Since the mean radius of the ψ' is much larger than that of the J/ψ, the equally strong suppression indicates that the final state interactions occur with a state that is not an eigenstate of the $(c\bar{c})$ system but rather with a common precursor of the J/ψ and ψ'. This does not come as a surprise, because the center of mass of the produced $(c\bar{c})$ pair moves rapidly with respect to the target nucleus. As viewed from the rest frame of the $(c\bar{c})$ pair, the nucleus is highly Lorentz contracted and thus the $(c\bar{c})$ pair leaves the target nucleus before the components corresponding to different quantum mechanical eigenstates of the charmonium system have had time to decohere.

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[1]This argument is not strictly correct, because heavy quarks are predominantly formed by gluon fusion or gluon scattering. Gluons might be more susceptible to initial state interactions than quarks and antiquarks. There is, indeed, evidence for enhanced initial state scattering of gluons from the p_T-spectra of heavy quark states produced in $p+\bar{p}$ reactions, giving rise to a broader "intrinsic" transverse momentum spread of gluons [4].
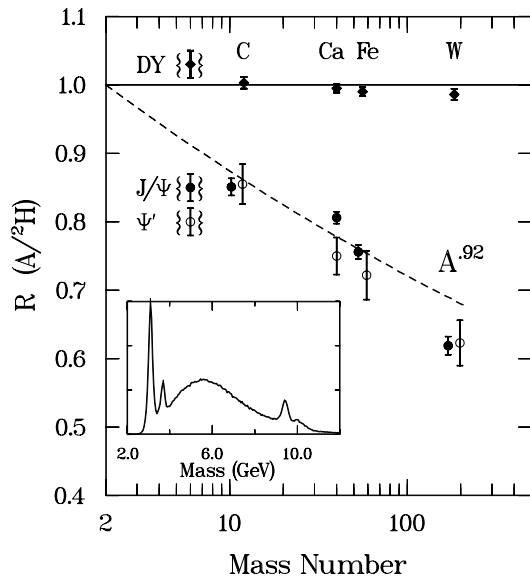
Figure 1: Ratio of nuclear cross section compared to deuterium for Drell-Yan pairs (DY), J/ψ and ψ' production. Note that the nuclear suppression for J/ψ and ψ' is the same. Data from experiment E772 at Fermilab [5].

The influence of the nuclear target on the decoherence process has been studied extensively in different theoretical approaches. However, the size of the nuclear suppression remained a great puzzle until quite recently. The problem is that the coherent superposition of eigenstates which is initially produced is characterized by a small geometric size of order $m_c^{-1} \approx 0.1$ fm. On the other hand, the sublinear rise with A of the charmonium production cross sections requires a final state interaction cross section $\sigma_{c\bar{c}N}$ of the order of several millibarn. A comprehensive analysis [6] of all available data on charmonium production in $p + A$ collisions within the Glauber model yields a value for the absorption cross section of $\sigma_{c\bar{c}N}^{(\mathrm{abs})} = (7.3 \pm 0.6)$ mb.

2.2   The Color Octet Model

The resolution of this apparent paradox can be achieved if one assumes that the $(c\bar{c})$ pair is originally produced in a color octet state [7] which strongly interacts with its environment. The color octet model was originally motivated by unexpectedly large cross sections for charmonium production observed at the Tevatron. Perturbative calculations fell short of the measured cross section for J/ψ and ψ' production at high p_T by at least one order of magnitude. The observation [8] that the data obtained by the CDF collaboration can be explained if high-p_T charmonia are predominantly formed by gluon fragmentation, led to the hypothesis that the charmonium wavefunctions contain a significant component where the

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($c\bar{c}$) pair is in a color octet state accompanied by a soft (valence) gluon:

$$|\psi(n^3\text{S}_1)\rangle = O(1)|c\bar{c}[^3\text{S}_1^{(1)}]\rangle + O(v_c)|c\bar{c}[^3P_J^{(8)}]g\rangle + O(v_c^2)|c\bar{c}[^3S_1^{(8)}]gg\rangle + \dots. \qquad (1)$$

Here v_c denotes the velocity of the bound charm quarks. This decomposition of the charmonium wavefunction can be unambiguously defined in the framework of the nonrelativistic expansion of QCD (NRQCD) developed by Lepage and collaborators [9]. The higher Fock space components are suppressed by powers of the velocity v_Q of the heavy quark. In the heavy quark limit $v_Q \approx \frac{1}{2}\alpha_s(m_Q) \ll 1$, producing a well ordered expansion into components with an increasing number of valence gluons. The reason for the suppression of higher Fock space components is that "physical" gluons couple to the color current of the heavy quarks which is proportional to their velocity $v_Q = p_Q/m_Q$.

In the limit $\alpha_s \ll 1$, where the structure of the heavy quark bound states is essentially Coulombic, the momentum scale for the production of the heavy quark pair, $2m_Q$, is much larger than the momentum scale associated with the heavy quark bound state, $p_Q(n^3S_1) = \alpha_s m_Q/2n$. In the case of the charmonium ground state, $p_c \approx 1$ GeV compared with $2m_c \approx 3$ GeV, which still provides for a reasonable separation of scales. The total S-matrix element for the production of a $c\bar{c}[n^3S_1]$ state may therefore be factorized into a perturbative matrix element for the elementary production process, such as gluon fusion ($gg \to c\bar{c}$) or gluon fragmentation ($g^* \to c\bar{c}$), and a nonperturbative matrix element $\mathcal{M}(c\bar{c} \to \psi(n^3S_1))$ describing the evolution of the elementary ($c\bar{c}$) pair into the strong interaction eigenstate (1), as illustrated in Fig. 2. There are different matrix elements, $\mathcal{M}_1$ and $\mathcal{M}_8$, depending on the color channel in which the ($c\bar{c}$) pair is produced.
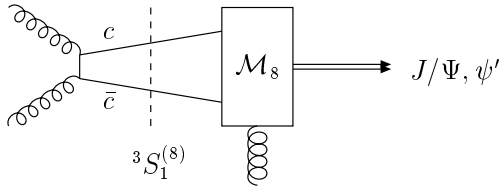

Figure 2: Schematic representation of the color-octet formation mechanism for the charmonium states J/ψ and ψ'. The intermediate state $^3S_1^{(8)}$ is a color octet which combines with a soft gluon to form the color-singlet charmonium state.

Assuming that charmonium production at high p_T at the Tevatron is dominated by the color octet channel, the matrix elements $\mathcal{M}_8$ for the formation of various eigenstates of the charmonium system can be deduced empirically [10, 11]. Their values are found to scale in accordance with the counting rules, the powers of v_c, predicted by (1). Although the

color octet model yields an excellent description of the measured differential cross sections, one has to caution that crucial predictions, such as the strong transverse polarization of J/ψ produced at high p_T, have not yet been confirmed. It should also be noted that the confirmation of this prediction would not automatically prove that the color octet mechanism also dominates J/ψ production at small p_T, where the cross section is concentrated in nuclear collisions.

An important theoretical problem is how quickly the color octet state neutralizes its color charge. Since the neutralization mechanism is, by assumption, a nonperturbative soft process, it is not rigorously calculable, but some qualitative estimates are possible. Since the force between the $(c\bar{c})$ pair in a color octet state is repulsive, the formation of a bound state requires that the relative $(c\bar{c})$ wavefunction must be a color singlet at distances comparable to the average radius of the charmonium state. More precisely, the relative magnitude of the octet and singlet components as a function of separation r should behave as

$$\frac{|\psi_{c\bar{c}}^{(8)}(r)|^2}{|\psi_{c\bar{c}}^{(1)}(r)|^2} \sim \exp\left[-2\int_0^r dr'\sqrt{m_c|E_{c\bar{c}} - V_8(r')|}\right]. \tag{2}$$

A crude estimate is obtained by replacing $|E_{c\bar{c}} - V_8|$ with the binding energy of the charmonium state, $E_B(n^3S_1)$, yielding the neutralization time

$$\tau_{8\rightarrow 1} \approx \left(2v_c\sqrt{m_c E_B(n^3S_1)}\right)^{-1} \approx 0.3n \text{ fm}/c. \tag{3}$$

Clearly, a more quantitative estimate within the framework of a coupled-channel potential model for charmonium would be very useful. In the laboratory frame, the octet-to-singlet transition time is Lorentz dilated due to the rapid motion of the $(c\bar{c})$ pair with respect to the target nucleus. Since the average traversed length of nuclear matter is of the order of 5 fm for a heavy nuclear target, the precise value of $\tau_{8\rightarrow 1}$ becomes important for charmonium states produced at midrapidity. However, in first approximation, it is reasonable to expect that the color neutralization occurs predominantly outside the target nucleus in $p + A$ experiments performed at several hundred GeV/c beam energy.

2.3   Nucleus-Nucleus Collisions

The production of J/ψ and ψ' in nucleus-nucleus collisions has been studied in the experiments NA38 and NA50 at CERN. NA38 investigated p + W, p + U, ^{16}O + U, and ^{32}S + U collisions at 200 GeV/nucleon; NA50 took data for ^{208}Pb + Pb collisions at 158 GeV/nucleon.
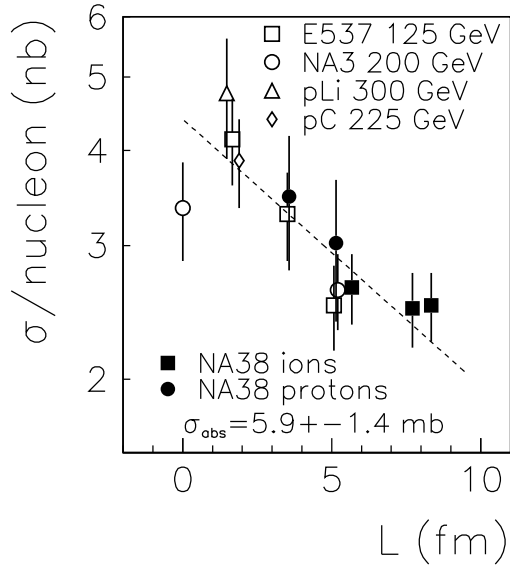
Figure 3: Effective nuclear absorption cross section in the Glauber approximation for J/ψ deduced from the NA38 data. The heavy ion data (black squares) fall on the same trend as the data from proton induced reactions.
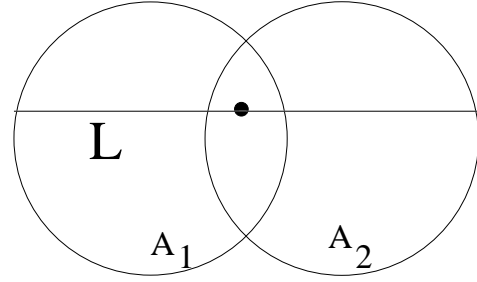

Figure 4: Illustration of the definition of L, the equivalent thickness of nuclear matter sweeping over the location of the created $(c\bar{c})$ pair.

The results of experiment NA38 showed that the production of J/ψ in ^{16}O- and ^{32}S-induced reactions follows the nuclear mass dependence expected from $p + A$ collisions see Figure (3).

The total production cross section $\sigma_{J/\psi}$ grows like[2]

$$\sigma_{J/\psi} = \sigma_{J/\psi}^{(pp)} \cdot A_1 A_2 \exp\left(-\sigma_{c\bar{c}N}^{(abs)} \rho_N L(A_1, A_2)\right) \equiv \sigma_{J/\psi}^{(pp)} A_1 A_2 S_N, \tag{4}$$

where $L(A_1, A_2)$ is the average length of nuclear matter traversed by the $(c\bar{c})$ pair after its formation at some moment during the collision. The definition of L is illustrated in Fig. 4.

In the framework of Glauber theory, the nuclear suppression factor S_N (and hence $L(A_1, A_2)$) is defined as

$$S_N = \int d^2b \, d^2b' dz dz' \rho_{A_1}(\vec{b}', z) \rho_{A_2}(\vec{b} - \vec{b}', z') T_{A_1}(z, \vec{b}') T_{A_2}(z', \vec{b} - \vec{b}') \tag{5}$$

with the nuclear profile function

$$T_A(z, \vec{b}) = \exp\left(-(A - 1)\sigma_{c\bar{c}N}^{(abs)} \int_z^\infty dz' \rho_A(\vec{b}, z')\right). \tag{6}$$

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[2]Note that $\rho_N L$ is a Lorentz invariant expression, which counts the area density of nucleons interacting with the $(c\bar{c})$ pair after its formation.

Here $\rho_A(\vec{b}, z)$ is the nuclear density $\sigma_{c\bar{c}N}^{(\text{abs})}$ distribution, normalized to unity so that $S_N = 1$ for $p + p$ collisions.

The fact that nuclear production of J/ψ appears to follow the Glauber model prediction (4) may be interpreted to mean that nothing "abnormal" occurs in nuclear reactions up to projectile mass $A_1 = 32$. This argument was first presented in this form by Capella, et al. [12] and by Gerschel and Hüfner [13]. It is worth noting that their conclusion originally was quite controversial because the absorption cross section $\sigma_{c\bar{c}N}^{(\text{abs})}$ was expected to be much smaller than the required 6–7 mb. This state of affairs has been radically changed by the introduction of the color-octet model that provides a natural explanation for $\sigma_{c\bar{c}N}^{(\text{abs})}$.

The results obtained for ψ' production on the other hand, cannot be explained by (4). The ψ' cross section is found to be more suppressed in ^{32}S + U reactions than the J/ψ cross section. This is in marked contrast to $p + A$ collisions where the ratio $\sigma_{\psi'}/\sigma_{J/\psi}$ is independent of the target mass. The different behavior of the ψ' becomes even more evident when the production of charmonium states is studied as a function of the total transverse energy E_T carried by the collision fragments. Within uncertainties, E_T is a measure of the impact parameter of the nuclear reaction. Selected ranges of E_T, therefore, correspond to different values of the average target thickness $L(A_1, A_2)$ in (4). The relation between L and E_T can be modeled in the Glauber approximation or some other geometric collision model. The relative suppression of ψ' production compared with J/ψ production increases with growing L (or E_T), as shown in Fig. 5 . On the other hand, $\sigma_{J/\psi}$ in different E_T-windows follows exactly the Glauber model formula (4), if the E_T-dependence of L is taken into account.

The stronger suppression of the ψ' indicates the presence of a new suppression mechanism in S + U collisions that is absent in $p + A$ collisions. One candidate for this mechanism is absorption by "comovers", i.e. by secondary hadrons produced at about the same rapidity as the ψ'. Since they are slowly moving with respect to the ψ', there is no time dilation of the formation time of the ψ' for interactions with these secondary particles. We will return to this issue in section 4, when we discuss models for comover absorption.

In Pb + Pb collisions at 158 GeV/nucleon, as shown in Fig. 6, also the cross section for J/ψ production is found to be more strongly suppressed than predicted by (4). The amount of additional suppression increases with L (or E_T), similar to the effect observed for ψ' production in S + U collisions. For ψ' production in the Pb + Pb system, strong additional suppression is observed even in the lowest E_T range, as shown in Fig. 7. The suppression factor for the highest E_T (or L) window is only insignificantly larger than that found in the most central S + U collisions. Taken together, the S + U and Pb + Pb data
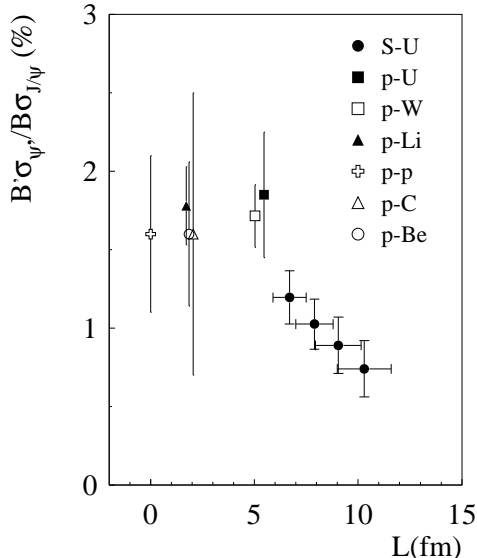
Figure 5: The cross section for ψ' production is suppressed relative to that for J/ψ production in S + U collisions. The suppression increases with L. Heavy ion data from experiment NA38.

indicate the onset of a new, "abnormal" suppression mechanism for charmonium production, first for the ψ' and later for the J/ψ, which appears to approach saturation for the ψ' in the most central Pb + Pb events.

The most obvious difference between $p + A$, S + U, and Pb + Pb interactions at SPS beam energies is the number of secondary particles. As already mentioned before, a significant fraction of the secondary hadrons are produced in the same region of rapidity as the charmonium. Is it possible to explain the "abnormal" suppression as absorption of the J/ψ and ψ' on these comoving hadrons? At first glance, this explanation seems eminently plausible. The absorption cross section of the ψ' and comovers should be much larger than that of the J/ψ, both because of the larger geometrical size of the ψ' and its lower energy threshold for dissociation. Therefore, the comover effect should set in much earlier for the ψ' than for the J/ψ, just as observed. The question is thus decidedly of a quantitative, rather than qualitative, nature:

1. Is it possible to explain the A- and E_T-dependence of the "anomolous" part of J/ψ and ψ'-suppression by a common set of parameters?

2. Do the required parameters, i.e. comover absorption cross sections, agree with theoretical estimates or values deduced from other relevant data?

The following sections address these two questions. We shall first review what theory can tell us about cross sections for the absorption of charmonium on hadrons made of light
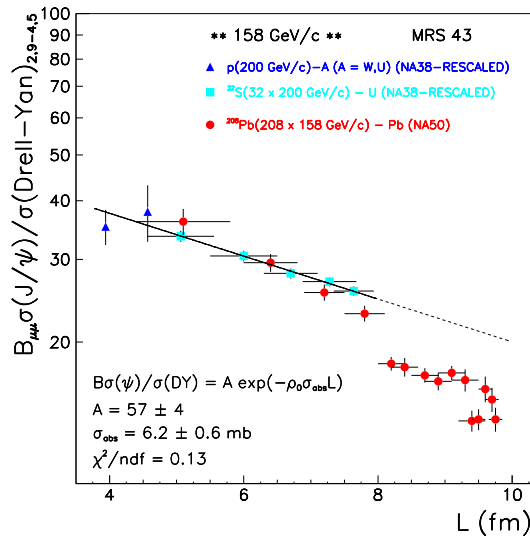
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Figure 6: Results from NA50 for the suppression of J/ψ production in Pb + Pb collisions as function of the Glauber parameter L. (From [14].)
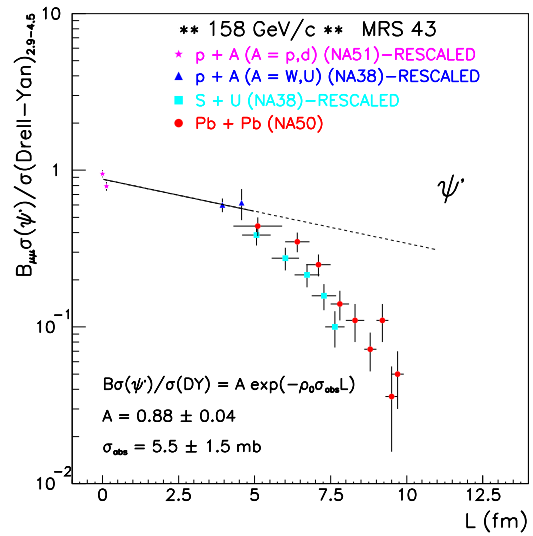

Figure 7: Results from NA50 for the suppression of ψ' production in Pb + Pb collisions as function of the Glauber parameter L. (From [14].)

quarks. We shall then turn to phenomenological models that aim at finding a consistent set of parameters describing the S + U and Pb + Pb data.

3 Charmonium Absorption by Comovers

3.1 Overview of Mechanisms

Most mechanisms that have been proposed to describe the absorption of J/ψ or ψ' on comoving hadronic matter fall into three categories:[3]

1. *Deconfinement:* If the comoving matter density is sufficiently high, the matter may be in the form of a quark-gluon plasma. In this case the charmonium states cannot form at all, as first argued by Matsui and Satz, because the attractive color force between the ($c\bar{c}$) pair is screened by the plasma. The color-singlet ($c\bar{c}$) potential then has the form

$$V_{cc}(r) = \frac{4}{3}\alpha_s \frac{e^{-\mu r}}{r}, \tag{7}$$

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[3]A comprehensive description of QCD based theoretical approaches to J/ψ interactions in hadronic matter can be found in [15].

where both α_s and μ are functions of the density of the plasma. For a thermalized plasma with no net baryon excess, finite temperature perturbation theory predicts

$$\mu^2 = \left(1 + \frac{1}{6}N_{\mathrm{f}}\right) g^2 T^2 \tag{8}$$

where $N_{\mathrm{f}} = 3$ is the number of light quark flavors. Monte-Carlo simulations of the SU(3) lattice gauge theory at finite temperature [16, 17] show that the inverse color screening length μ in the range $T_{\mathrm{c}} \leq T \leq 3T_{\mathrm{c}}$ can be parametrized as $\mu \approx 2.2T$. For every bound state of a heavy quark-antiquark pair then exists a critical temperature T_{D}, above which the state disappears. Detailed studies [2] have revealed that $T_{\mathrm{D}}(\psi') \approx T_{\mathrm{D}}(\chi_c) \approx T_{\mathrm{c}}$, $T_{\mathrm{D}}(J/\psi) \approx 1.2T_{\mathrm{c}}$, and $T_{\mathrm{D}}(\Upsilon) = 2T_{\mathrm{c}}$.

2. *Thermal dissociation:* If the $(c\bar{c})$ pair is immersed in a thermal hadronic gas, its internal modes will eventually also become thermally occupied. All those states above the $D\bar{D}$ threshold will dissociate. As the temperature rises, the thermal fraction of dissociated states increases until, eventually, the bound states make up a negligible contribution. In condensed matter physics, the temperature above which almost no bound states remain is called the Mott temperature. In this picture, the J/ψ and ψ' could essentially disappear even if quarks and gluons remain confined at any temperature. However, the critical question is how fast the internal states of a heavy quark pair become thermalized by interactions with a hadronic medium. This is an issue of kinetics, rather than thermodynamics, and will be addressed in the next subsection.

3. *Pre-thermal dissociation:* In this scenario, charmonium states could be excited above the dissociation $(D\bar{D})$ threshold by interactions with a medium of comovers that had no time to equilibrate. Two popular examples are (i) a pre-thermal parton plasma [18, 19], (ii) a system of color flux tubes [20, 21].

The first scenario, a pre-equilibrium parton plasma, obviously only makes sense if the dissociation process is faster than the thermalization of the plasma: $\tau_{\mathrm{dis}} \leq \tau_{\mathrm{eq.}}$. This seems an unlikely case, because parton thermalization is governed by the cross section σ_{gg}. But in the perturbative realm of QCD, $\sigma_{\mathrm{gg}} \approx \frac{9}{4}\sigma_{\mathrm{gc}}$, hence the plasma should equilibrate faster than a $(c\bar{c})$ state can be dissociated.

Nonetheless, quantitative model studies of the competition between equilibration of partons and J/ψ dissociation would be useful. These could be carried out at two levels. One could study the dissociation of a fully formed charmonium state by gluon impact in

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the environment of a gluon distribution that itself is equilibrating. Or one could study the influence of a prethermal gluon bath on the conversion process from color octet $(c\bar{c})$ to the color singlet charmonium state. The latter would require some modeling of the nonperturbative matrix element $\mathcal{M}_\mathrm{g}$, but a semi-quantitative treatment of this mechanism should be possible.

The second scenario, dissociation by color flux tubes, again requires that the decay rate of flux tubes by light quark pair production, $\tau_{q\bar{q}}$, is slower than that of the dissociation rate of the charmonium state. This also appears unlikely, because QCD flux tubes are known to break rapidly ($\tau_{q\bar{q}} < 1$ fm/c), while a $(c\bar{c})$ pair cannot be pulled apart too rapidly by virtue of the large mass of the c-quark. Again, simple model studies of the competition between these two processes would be easily carried out, and the dependence on heavy quark mass investigated. In this scenario, too, one can study the possible influence of the presence of coherent gluonic fields on the color octet-to-singlet conversion process.

3.2 Thermal Dissociation Kinetics

Since, in general, the cross sections σ_hh among light hadrons are much larger than those between light hadrons and charmonium, $\sigma_{\mathrm{h}\psi}$, it also makes sense to restrict the study of dissociation of charmonium by light hadrons to a thermal hadronic environment.[4] As we argued before, the same consideration applies to a partonic environment above T_c.

There are two possible approaches to the dissociation problem: one at the quark level where the large mass of the c-quark is used to separate perturbative from nonperturbative aspects of the problem, and another one that makes use of effective hadronic interactions. In this section, we will consider the constituent $(c\bar{c})$-approach. An effective hadronic Lagrangian approach will be discussed in the following section.

As first explored in detail by Peskin [22] and Bhanot and Peskin [23], interactions between heavy quark bound states and light hadrons can be described perturbatively, if the heavy quark mass m_Q is large enough. The reason is that the characteristic momentum scale of the heavy quark bound state is of order $\alpha_\mathrm{s} m_\mathrm{Q}$ which is large. As a result, only gluons with momenta larger than $\alpha_\mathrm{s} m_\mathrm{Q}$ can resolve the internal color structure of the $(Q\bar{Q})$ bound state and interact with it. To gluons with much smaller momenta, the $(Q\bar{Q})$ state appears as an inert color singlet. The small size of the $(Q\bar{Q})$ state then allows for a systematic multipole

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[4]The validity of this assumption can be checked with the help of hadronic cascade models.

expansion of its interaction with external glue fields, where the dipole interaction dominates at long range.

The potential problem with this approach is that the c-quark is not quite heavy enough. For instance, the binding energy of the J/ψ is about 650 MeV, which is neither small compared to the reduced mass $\frac{1}{2}m_c \approx 700$ MeV of the charmonium system, nor large on the scale of light hadron masses. Gluons with momenta of order 1 GeV/c or even slightly less can effectively resolve the internal color structure of the J/ψ, but it is not clear that they belong in the perturbative domain of QCD.
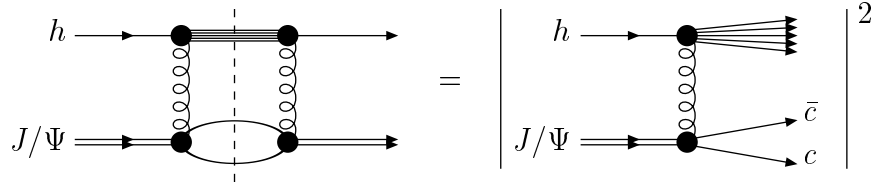

Figure 8: Forward scattering between a light hadron and a J/ψ is, in leading order, described as two-gluon exchange. The total absorption cross section is related to the imaginary part of this forward scattering amplitude.

By virtue of the optical theorem, the absorptive cross section between a light hadron h and a charmonium J/ψ can be expressed as the imaginary part of the forward scattering amplitude where the interaction is dominated at long distances by two-gluon exchange (see Figure 8). In this picture, light hadrons interact with the J/ψ only via their glue content. Kharzeev and Satz [24] applied the Bhanot-Peskin formalism to $\pi - J/\psi$ scattering. Using the operator product expansion to separate long-distance physics from the (assumed) perturbatively calculable structure of the $(c\bar{c})$ system, they found that the absorptive cross section is proportional to the part of the gluon structure function $G_h(x)$ of the interacting hadron that is sufficiently energetic to dissociate the J/ψ. Since the gluon structure functions of light hadrons typically are soft, i.e. $G_h(x) \xrightarrow{x\to1} (1-x)^5$, only highly energetic hadrons are capable of exciting a J/ψ above the dissociation threshold. The peak cross section[5] for J/ψ-dissociation by a gluon in the dipole approximation is about 3 mb at a gluon momentum around 1 GeV/c. Due to the softness of the gluon structure function, the pion momentum must reach 5 GeV/c before attaining an absorption cross section in excess of 1 mb. The

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[5]The fact that $\sigma_{g\psi}^{(abs)}$ peaks and then falls off is presumably an artifact of the dipole approximation. If higher orders were retained, the cross section should approach a constant value at high energy, $\sigma_{g\psi}^{abs} \to 2\sigma_{gc}$.

thermally averaged cross section, in this framework, remains less than 0.1 mb for a pion gas within any realistic temperature range.
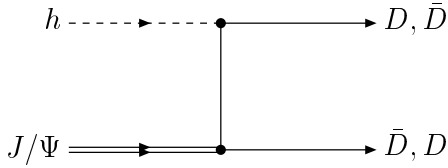

Figure 9: Inelastic scattering of a light hadron on a J/ψ, yielding two D-mesons in the final state.
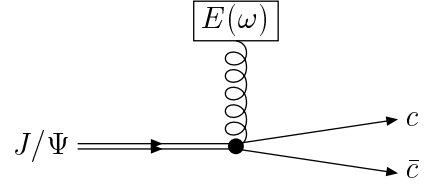
Figure 10: The J/ψ state is dissociated into a $(c\bar{c})$ octet pair by absorption of an energetic gluon from its environment.

The calculation of [24] is clearly incomplete, because pions are not very efficient at dissociating the J/ψ. Heavier hadrons, such as the ρ-meson, can make significant contributions although they are much less abundant than pions. Moreover, the calculation [24] does not account for the energy balance between initial and final states, i.e. that the total mass of the colliding hadron is available to convert the J/ψ state into a pair of D-mesons (see Fig. 9). This effect is especially important for heavier mesons, such as the ρ.

Repeating the calculation [24] in a somewhat more general form, one finds that the dissociation rate of J/ψ-mesons in a (thermal) hadronic medium is given by (see Appendix A for details)

$$\Gamma_{\mathrm{dis}} \approx \frac{1}{\Delta t} \int_{\omega_{\mathrm{th}}}^{\infty} d\omega \langle |E(\omega)|^2 \rangle \, \pi \alpha_s a^2, \tag{9}$$

where a is the Bohr radius of the J/ψ and $\frac{1}{\Delta t}\langle |E(\omega)|^2 \rangle$ denotes the time-average fluctuation density of color (electric) fields with frequency ω in the hadronic environment, as illustrated in Fig. 10. The lower limit of the integral accounts for the dissociation energy threshold, and the factor $(\alpha_s a^2)$ arises from the coupling to the color dipole moment of the $(c\bar{c})$ pair in the J/ψ. Equation (9) is applicable to any kind of hadronic medium, whether thermal or not. For example, one could estimate the color field fluctuations in the framework of a random field model as proposed by Hüfner et al. [25]. For a thermalized medium, QCD sum rules or lattice gauge theory could be applied. Because some heavy quarkonium states, especially the Υ do not disappear right at the critical temperature, the color field fluctuation spectral density is also an important quantity to know in the deconfined phase, where it could be evaluated either in perturbation theory or by lattice simulations.
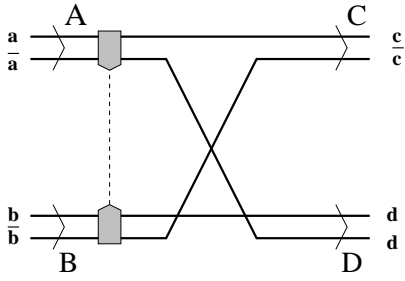
Figure 11: In the framework of the constituent quark model, the absorption of a J/ψ on a light quark meson is viewed as quark exchange reaction. At lowest order, four one-gluon exchange diagrams contribute to the amplitude. (From Martins et al. [27].)

The real limitations of (9) lie in the fact that the J/ψ is not truly a perturbative bound state but probes also the confining part of the $c\bar{c}$-potential. It is therefore of interest to investigate the dissociation rate also in the framework of other approaches that model quark confinement. One such model is the constituent quark model of Isgur and Karl [26]. In this model the reaction

$$\pi + J/\psi \rightarrow D^* + \bar{D} \quad (D + \bar{D}^*) \tag{10}$$

is viewed as quark exchange where a c-quark and a light quark change sides. The cross section for this reaction was calculated by Martins, Blaschke, and Quack [27] in the first Born approximation (see the diagrammatic representation in Fig. 11). Including $D^*\bar{D}$, $D\bar{D}^*$, and $D^*\bar{D}^*$ final states, the cross section peaks around 1 GeV kinetic energy (c.m.) at a value of 15 mb (see Fig. 11). The magnitude of this cross section is solely due to the action of the confining interaction between the quarks, which is modeled as a "color-blind" attractive interaction between the quarks which has a Gaussian momentum dependence. Because this interaction is taken as attractive independent of the color quantum numbers of the affected quark pair, it does not cancel for the interaction of a light quark with a point-like $(c\bar{c})$ pair in a color-singlet state, in contrast to the one-gluon exchange interaction. Moreover, the magnitude of the obtained cross section invalidates the use of the Born approximation, unless a number of different partial waves contribute. This is impossible to ascertain from [27], because no partial wave decomposition of the dissociation cross section is presented there.

Since the results obtained within the constituent quark model differ by orders of magnitude from those obtained in the Bhanot-Peskin approach, it is useful to consider the absorption process also in an entirely different framework. This will be done in the next subsection.
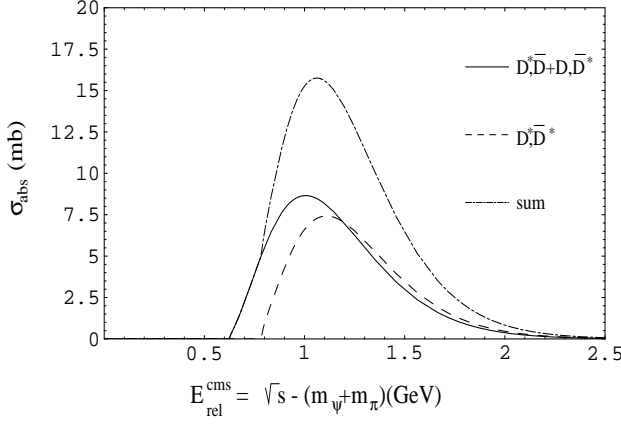
Figure 12: Energy dependence of the J/ψ absorption on pions for the final states DD^* and D^*D^*. (From the preprint version of Martins, et al. [27]. Note that the figure in the printed version differs slightly from the one reproduced here.)

3.3 Hadronic Dissociation: Effective Theory

The most abundant mesons in a hot hadronic gas are π, K, and ρ. These can induce the following dissociation processes when encountering a J/ψ particle:

$$\pi + J/\psi \to D + \bar{D}^*, \quad \bar{D} + D^* \tag{11}$$

$$K + J/\psi \to D_s + \bar{D}^*, \quad \bar{D}_s + D^*, \quad D + \bar{D}_s^*, \quad \bar{D} + D_s^* \tag{12}$$

$$\rho + J/\psi \to D + \bar{D}, \quad D^* + \bar{D}^* \tag{13}$$

The kinematic thresholds for these reactions are

(11)   640 MeV,   (12)   385 MeV,   (13) -135 MeV (+155 MeV).

The reaction $\rho + J/\psi \to D + \bar{D}$ has the lowest total invariant mass threshold and is, in fact, exothermic. At thermal equilibrium it is only suppressed due to the relatively high mass of the ρ-meson which causes ρ-mesons to be less abundant than pions in a thermal hadron gas.

From a microscopic point of view, all reactions listed above can be understood as quark exchanges, where the J/ψ transmits a charm quark to the light meson and picks up a light (u, d, or s) quark. Since similar reactions among light hadrons typically occur with large cross sections, one can expect that the reactions (11,12,13) also proceed with significant strengths above their respective kinematic thresholds. This is, indeed, the result obtained by Martins, Blaschke, and Quack [27] which was discussed at the end of the previous subsection.

Here, we follow a different approach [28]. In the effective meson theory, the exchange of a ($c\bar{q}$) or ($\bar{c}q$) pair, where $q, \bar{q}$ stands for any light quark, can be described as the exchange of

16

a D, D_s, D^*, or D_s^* meson between the J/ψ and the incident light meson (see Fig. 13). Near the kinematic threshold, where only a small number of final states are open, the description in terms of an effective meson exchange is expected to be valid.
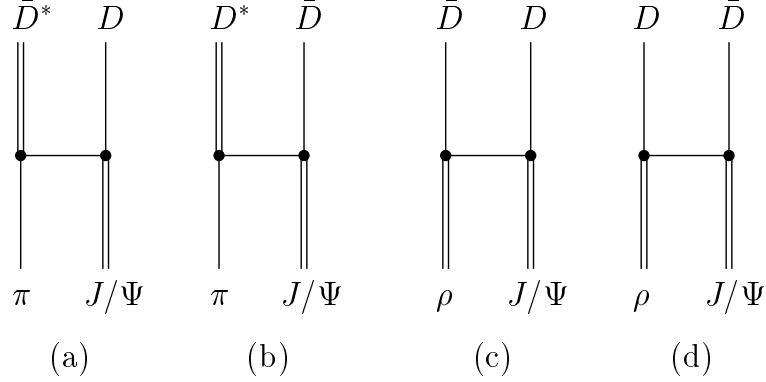

Figure 13: Lowest order Feynman diagrams contributing to the charm exchange reactions (11) and (13). Single lines represent pseudoscalar mesons; double lines denote vector mesons.

In order to calculate the various Feynman diagrams for the reactions (11,12,13) we need to construct the effective three-meson vertices. We do so by invoking a strongly broken U(4) flavor symmetry with the vector mesons playing the role of quasi-gauge bosons. Denoting the 16-plet of pseudoscalar mesons $(\pi, \eta, \eta', K, D, D_s, \eta_c)$ by $\Phi = \phi_i T_i$, where the T_i are the U(4) generators, and the vector meson 16-plet $(\omega, \rho, \phi, K^*, D^*, D_s^*, \psi)$ by $\mathcal{V}^\mu = V_i^\mu T_i$, the free meson Lagrangian reads

$$\mathcal{L}_0 = \text{tr}(\partial^\mu \Phi^\dagger \, \partial_\mu \Phi) - \text{tr}(\partial^\mu \mathcal{V}^{\dagger,\nu})(\partial_\mu \mathcal{V}_\nu - \partial_\nu \mathcal{V}_\mu) - \text{tr}(\Phi^\dagger M_\text{P} \Phi) + \frac{1}{2} \text{tr}(\mathcal{V}^{\mu\dagger} M_\text{V} \mathcal{V}_\mu). \qquad (14)$$

Here M_P and M_V denote the mass matrices for the pseudoscalar and vector mesons, respectively. Because of the heavy mass of the charm quark, M_P and M_V break the U(4) symmetry strongly down to U(3), the mass of the strange quark introduces a weaker breaking to U(2), and the axial anomaly further breaks the U(2) symmetry to SU(2) in the case of the pseudoscalar mesons. All these symmetry breakings are embodied in the physical mass matrices. It is convenient, in the following, to work with the mass eigenstates.

The meson couplings are obtained by replacing the space-time derivatives ∂_μ by the "gauge covariant" derivatives

$$D_\mu = \partial_\mu - ig\mathcal{V}_\mu. \qquad (15)$$

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In first order in the coupling constant g this procedure leads to the following interactions

$$
\begin{aligned}
\mathcal{L}_{\text{int}} &= ig \operatorname{tr}\left(\Phi^\dagger \mathcal{V}^{\mu\dagger}\partial_\mu\Phi - \partial^\mu\Phi^\dagger\mathcal{V}_\mu\Phi\right) \\
&\quad + ig \operatorname{tr}\left(\partial^\mu\mathcal{V}^{\dagger\nu}[\mathcal{V}_\mu, \mathcal{V}_\nu] - [\mathcal{V}^{\dagger\mu}, \mathcal{V}^{\dagger\nu}]\partial_\mu\mathcal{V}_\nu\right).
\end{aligned}
\tag{16}
$$

If the U(4) flavor symmetry were exact, we would expect all couplings given by the same constant g. In view of the significant breaking of the flavor symmetry we anticipate that the effective coupling constants for different 3-meson vertices will have different values. We will see below to what extent this is true.

In order to describe π- and ρ- induced J/ψ dissociation, we need the following vertices: ψDD, ψD^*D^*, πDD^*, ρDD, and ρD^*D^*. From (16) we derive the following interactions:

$$
\begin{aligned}
\mathcal{L}_{\psi DD} &= ig_{\psi DD}\psi^\mu\left(\bar{D}\partial_\mu D - (\partial_\mu\bar{D})D\right) \\
\mathcal{L}_{\psi D^*D^*} &= -ig_{\psi D^*D^*}\psi^\mu\left(\bar{D}^{*\nu}\partial_\mu D_\nu^* - (\partial_\mu\bar{D}^{*\nu})D_\nu^*\right) \\
\mathcal{L}_{\pi DD^*} &= \frac{i}{2}g_{\pi DD^*}\left(\bar{D}\tau_i D^{*\mu}\partial_\mu\pi_i - \partial^\mu D\tau_i D_\mu^*\pi_i - \partial_\mu\pi_i\bar{D}^{*\mu}\tau_i D + \pi_i\bar{D}^{*\mu}\tau_i\partial_\mu D\right) \\
\mathcal{L}_{\rho DD} &= \frac{i}{2}g_{\rho DD}\rho_i^\mu\left(\bar{D}\tau_i\partial_\mu D - \partial_\mu\bar{D}\tau_i D\right) \\
\mathcal{L}_{\rho D^*D^*} &= -\frac{i}{2}g_{\rho D^*D^*}\rho_i^\mu\left(\bar{D}^{*\nu}\tau_i\partial_\mu D_\nu^* - (\partial_\mu\bar{D}^{*\nu})\tau_i D_\nu^*\right).
\end{aligned}
\tag{17}
$$

The coupling constants $g_{\psi DD}$, $g_{\psi D^*D^*}$, $g_{\rho DD}$ and $g_{\rho D^*D^*}$ can be derived from the D and D^* electric form factors in the framework of the vector meson dominance (VMD) model.

If γ_V denotes the photon-vector meson $\mathcal{V}$ mixing amplitude, the standard VMD analysis [29] yields the relations

$$
\gamma_\rho f_\rho = em_\rho^2, \qquad \gamma_\psi f_\psi = \frac{2}{3}em_\psi^2
\tag{18}
$$

where universality of the vector meson couplings is assumed:

$$
f_\rho = g_{\rho DD} = g_{\rho D^*D^*}, \qquad f_\psi = g_{\psi DD} = g_{\psi D^*D^*}.
\tag{19}
$$

The photon mixing amplitudes γ_V can be determined from the leptonic vector meson decay widths [29]:

$$
\Gamma_{Vee} = \frac{1}{3}\alpha\frac{\gamma_V^2}{m_V^3}.
\tag{20}
$$

Inserting the experimental numbers, we find

$$
f_\rho \approx 5.6, \qquad f_\psi \approx 7.7.
\tag{21}
$$

18

As is well known, different ways of deriving the value of the "universal" ρ-meson coupling f_ρ yield values differing by about 20%. The coupling constants (21) must therefore be considered to be given with an error of this order of magnitude.

Finally, the π-meson coupling between D and D^* can be obtained from the decay width of the D^*-meson: $D^* \to D\pi$. Unfortunately, only an upper bound for this decay rate is known at present, corresponding to $g_{\pi DD^*} < 15$. Theoretical estimates for this decay rate, based on QCD sum rules [30], yield

$$g_{\pi DD^*} \approx 8.8. \tag{22}$$

We adopt this value here.

At the hadronic level, the charm exchange reaction between the J/ψ and a light hadron can proceed either by exchange of a D- or a D^*-meson. Here we shall only consider the D-exchange reactions, because they have an acceptable high-energy behavior even in the absence of form factors for the vertices. The D^*-exchange reactions have cross sections that rise rapidly with energy due to the exchange of longitudinally polarized D^*-mesons. This growth characteristic of the exchange of massive vector bosons is obviously unphysical and will be cut off by the mesonic form factors long before it becomes significant. However, this would require that we introduce an unknown parameter (the characteristic momentum scale of the internal mesonic structure). In order to avoid this complication, we here neglect D^*-exchange cross sections. One final remark on D^*-exchange: Regge theory dictates that, in the high energy limit, the charm exchange reaction is dominated by the exchange of the D^*-trajectory. However, here we are not interested in charm exchange at high energies but near the kinematical threshold, because the relative motion of the hadrons is limited to thermal momenta.

Analytical expressions for the amplitudes and cross sections for the processes corresponding to the Feynman diagrams in Fig. 13 are given in Appendix B. The charm exchange cross sections are plotted in Fig. 14 as functions of the center-of-mass energy $\sqrt{s}$. The $\pi + J/\psi$ cross section (dashed line) starts at zero, because the reaction is endothermic, whereas the $\rho + J/\psi$ cross section (solid line) is finite at the threshold because of its exothermic nature. Of course, the reaction rate vanishes at threshold in both cases. Note that the cross sections for these two J/ψ absorption reactions are of similar magnitude over the energy range relevant to a thermal meson environment.

Figure 15 shows the pion and ρ-meson densities in an ideal thermal meson gas as a function of temperature. In the temperature range where a hadronic gas is likely to exist,
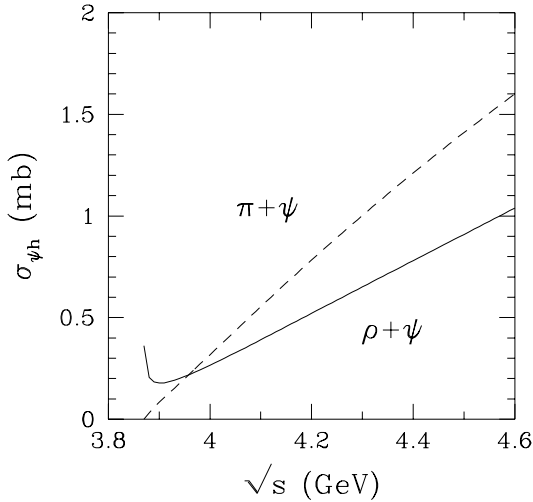
Figure 14: Cross sections for the charm exchange reactions described by the diagrams of Figure 13, as functions of c.m. energy. Dashed line: pions, solid line: ρ-mesons.
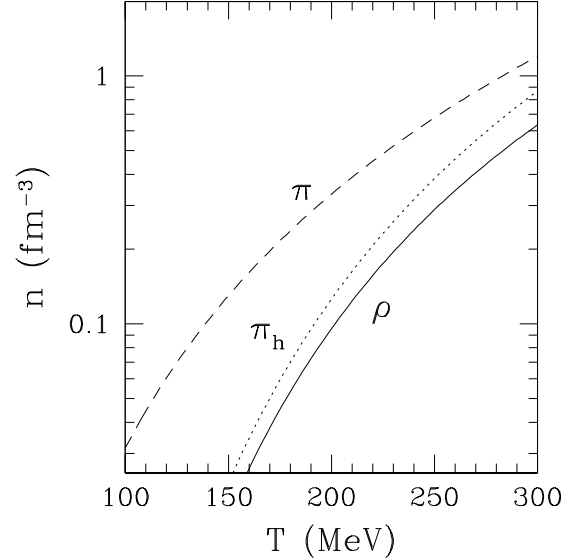
Figure 15: Thermal pion and ρ-meson densities in an ideal hadron gas as function of temperature. Dashed line: pions, solid line: ρ-mesons. The dotted line represents the density of pions above the dissociation threshold for J/ψ.

$T \leq 170$ MeV, pions are far more abundant than ρ-mesons, but most of these pions do not have sufficient energy to initiate the dissociation of a J/ψ. In fact, if one counts only those pions and ρ-mesons above the kinematical threshold for J/ψ-dissociation, π- and ρ-mesons are about equally abundant, or rather rare, because the effective density does not exceed 0.1 fm^{-3} even at $T = 200$ MeV, where the hadron picture of a thermal environment probably already fails. Figure 16 shows the J/ψ absorption rates in a thermal meson gas, as a function of temperature. Even at the (unrealistically) high temperature $T = 300$ MeV, the thermal dissociation rate is still so small that it corresponds to a lifetime around 10 fm/c. Thermal dissociation at $T \leq 200$ MeV is completely negligible on the time scale of the lifetime of a hot hadronic gas state in nuclear collisions.

One can ask the question whether a form factor should be included in the meson vertices used to evaluate the Feynman diagrams in Fig. 14. In principle, the concept of vector meson dominance assumes that the off-shell behavior of the vector meson propagators describes the form factor of the pseudoscalar mesons correctly, when the pseudoscalar mesons are on-shell. In the diagrams of Fig. 14, however, the exchanged D-meson is off-shell by a considerable
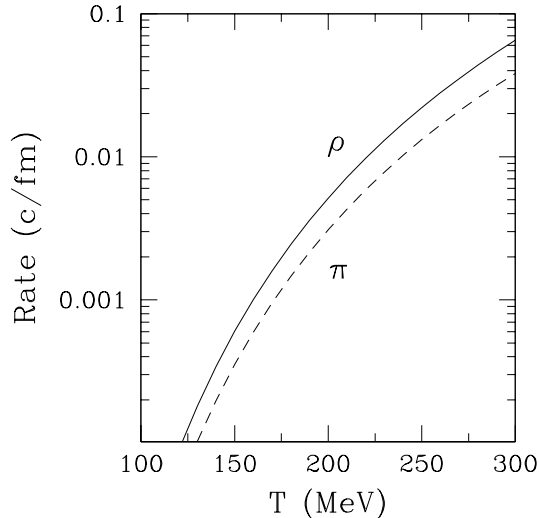
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Figure 16: Thermal J/ψ absorption rates as function of temperature. The pion and ρ-meson rates are shown separately.
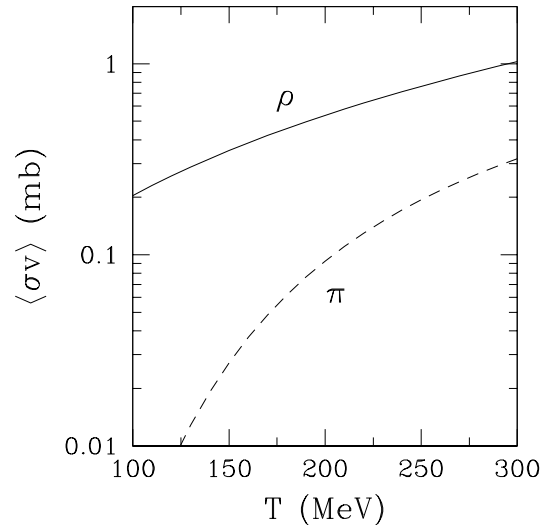
Figure 17: Thermally averaged J/ψ absorption cross sections as function of temperature. The pion and ρ-meson rates are shown separately.

amount, whereas the vector mesons all appear on-shell either in the initial or final state. It is unclear whether an additional form factor is needed in this situation. In any case, our result represents an upper limit, because an additional form factor would reduce the charm exchange rates. (E.g. a Gaussian form factor $\exp(-Q^2/Q_0^2)$ with $Q_0 = 1.5$ GeV would further reduce the absorption rates by slightly more than one order of magnitude.)

As we noted above, our analysis also remains incomplete because of the neglect of D^*-exchange reactions. It would be most interesting to consider also these reactions in the framework of a complete effective Lagrangian describing the interactions of pseudoscalar and vector mesons with quark content $(q\bar{q})$, $(Q\bar{q})$, $(q\bar{Q})$, and $(Q\bar{Q})$, where q stands for any light quark flavor and Q denotes a heavy quark. Such a Lagrangian embodying chiral symmetry for the light quarks has recently been given by Chan [31]. Loop corrections to the tree diagrams would also permit the study of form factor effects in this approach.

4   Models versus the Data

4.1 Glauber Theory

Probably the most conservative approach to calculate nuclear effects of suppression of charmonium production is based on the general framework of Glauber theory. This was, in fact, the starting point of the early analysis of Gerschel and Hüfner [13], and it forms the basis of more recent analyses of the available $p + A$ and $A + A$ data [6, 32]. In the Glauber formalism the suppression factor S_N due to absorption on nucleons is given by (5). The effect of absorption by comoving secondary hadrons is included by multiplying the integrand in (5) by an additional exponential absorption factor:

$$S = \int d^2b \, d^2b' \, \frac{dS_N}{d^2b \, d^2b'} \, T_{\rm co}(b, b') \tag{23}$$

with

$$T_{\rm co}(b, b') = \exp\left[-\sigma_{\rm co}^{\rm eff} \int_{\tau_0}^{\tau_f} \frac{dN_{\rm co}(b, b')}{dy} d\tau \right]. \tag{24}$$

Here $dN_{\rm co}(b, b')/dy$ is the comover density per unit of rapidity at impact parameter $\vec{b}, \vec{b'}$ with respect to the two nuclear centers, and $\sigma_{\rm co}^{\rm eff} = \langle \sigma_{\rm h\psi}^{\rm (abs)} v \rangle$ is the effective J/ψ absorption cross section by comovers. The effect of quark deconfinement (complete absorption) can be easily described in this framework by setting $\sigma_{\rm co}^{\rm eff} = \infty$ for $dN_{\rm co}/dy > dN_{\rm co}^{\rm crit}/dy$. τ_0 corresponds to the formation time of the dense comover gas (typically taken as $\tau_0 = 1 - 2$ fm/c), and τ_f denotes the time of hadronic freeze-out.

It is reasonable to assume that the comover density is proportional to the transverse energy dE_T/dy produced in a nucleus-nucleus collision at impact parameter b. A detailed analysis of dE_T/dy measurements in S + U and Pb + Pb collisions [6] shows that they can be well described by assuming that the transverse energy production is proportional to the number of participant projectile and target nucleons, as calculated in a geometrical picture or within the Glauber model itself. This is in agreement with a large body of data indicating that the "wounded nucleon" model [34] provides a good description of global observables in nuclear collisions. Using Bjorken's formula for the comoving energy density [35]

$$\epsilon_0 = \frac{dE_T/dy}{\pi R^2 \tau_0}, \tag{25}$$

where πR^2 denotes the effective nuclear geometric cross section, and setting $\tau_0 = 1$ fm/c, the initial energy densities for different impact parameters in Pb + Pb and S + U collisions are found as shown in Fig. 18. Clearly, for most impact parameters the density achieved in Pb + Pb collisions exceeds that produced in even the most central S + U collisions.
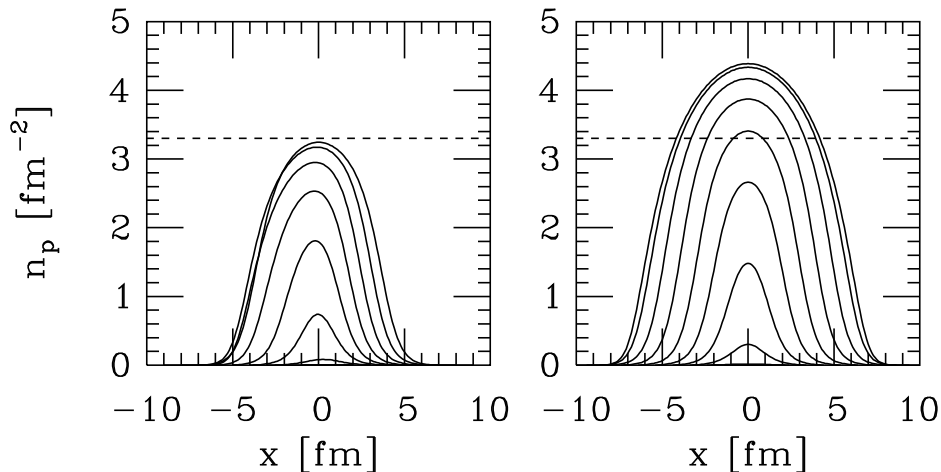
Figure 18: Initial energy densities according to the Bjorken formula (25) are shown as functions of impact parameter for (left) S + U and (right) Pb + Pb collisions. (From ref. [36].)

This observation opens the door [36] for an explanation of the anomalous J/ψ production observed in Pb + Pb collisions as the effect of deconfinement if a critical energy density $\epsilon_c \approx 2.75$ GeV/fm^3 is exceeded (more precisely, if $\epsilon_0 \tau_0$ exceeds 2.75 GeV/fm$^2 c$).

The earliest systematical analysis of suppression by comovers is due to Gavin and Vogt [37]. Using the parameters $\sigma_{c\bar{c}N}^{(\text{abs})} = 4.8$ mb, $\sigma_{\text{co}}^{\text{eff}} = 3.2$ mb, and $\tau_0 = 2$ fm/c that were originally derived from the S + U data [38] they find a good overall agreement with the measured suppression in Pb + Pb collisions [39]. However, it must be noted that, according to our present understanding, their value of $\sigma_{c\bar{c}N}^{(\text{abs})}$ may be too small to consistently explain the $p + A$ data [6], and $\sigma_{\text{co}}^{\text{eff}}$ is most likely much too large. The point here is that at $\tau_0 = 2$ fm/c the formation of a color singlet ($c\bar{c}$) state should be completed and, hence, the comover absorption should be of hadronic size ($\sigma_{\text{co}}^{\text{eff}} < 1$ mb) as discussed in the previous section.

On the basis of their Glauber model analysis of $p + A$ data on J/ψ suppression, Kharzeev et al. [6] argue that there is no room for absorption by comovers in S + U collisions (left part of Fig. 19). In order to explain the significant additional suppression observed in the Pb + Pb data (right part of Fig. 19), these authors must invoke a suppression mechanism that sets in abruptly and strongly if dE_T/dy, or rather the area density of wounded nucleons, exceeds a certain critical value.

This is also found in the analysis of the same data by Armesto and Capella [32], who
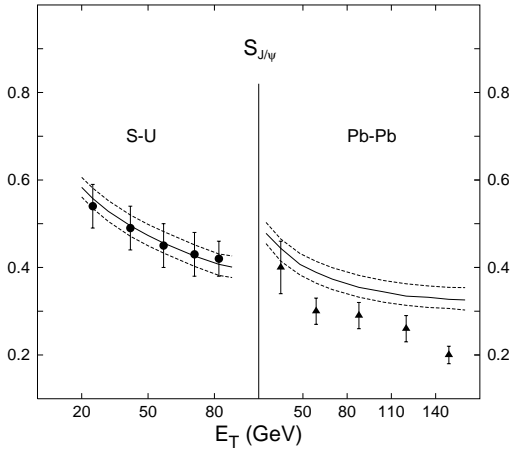
Figure 19: J/ψ suppression factor due to absorption on nucleons as obtained in the Glauber approximation with an absorption cross section of 7.3 ± 0.6 mb, in comparison with the NA38/50 data for the S + U and Pb + Pb systems (from [6]). The suppression in S + U is nicely explained, but the Glauber model fails in the Pb + Pb system. The excellent agreement for S + U appears to leave no room for comover absorption in this system.

obtain acceptable fits with the parameter choices $\sigma_{c\bar{c}N}^{(\mathrm{abs})} = 6.7(7.3)$ mb, $\sigma_{\mathrm{co}}^{\mathrm{eff}} = 0.55(1.0)$ mb, and critical area density of wounded nucleons of $dN^{\mathrm{crit}}/dy = 1.15(2.5)$ fm^{-2}.

What these two analyses demonstrate is that the present body of data on J/ψ production in $p + A$ collisions leave very little room for additional suppression effects in S + U collisions, at least within the framework of the Glauber model. In order to obtain sufficiently strong suppression effects in Pb + Pb, a strongly nonlinear dependence of the exponent of (24) on the comover density (here modeled as a threshold effect) is then required. Whereas Kharzeev et al.[6] argue that this indicates a qualitatively new mechanism (such as color deconfinement), the authors of [32] do not reach this conclusion.

The results for the hadronic absorption rate described above arguably point toward a failure of the Glauber model to describe the Pb + Pb data, at least if only hadronic interactions are considered.[6] Whether the reason for this failure is the emergence of a new suppression mechanism, or whether the A-dependence of the effective comover density is strongly underestimated by the Glauber approach, is impossible to tell from the existing data. The latter possibility has been recently addressed in [40], where it was pointed out that the comovers responsible for J/ψ suppression have to be "semihard" in order to resolve the color dipole structure of the J/ψ. A perturbative QCD picture then suggests that the area density of these comovers grows as $(A_1, A_2)^{1/3}$ in collisions at the SPS, much faster than the $(A_1^{1/3} + A_2^{1/3})$ growth predicted by the wounded nucleon model. Of course, such comovers would be better represented as partonic excitations, not hadronic ones. Indeed, a hybrid

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[6]Not all authors share this opinion. In particular, Gavin[33] and Capella[32] have maintained that absorption by hadronic or "pre-hadronic" comovers can describe all observations.

parton cascade model [41] predicts just such a dependence for the partonic components of secondary particles (see Fig. 20).
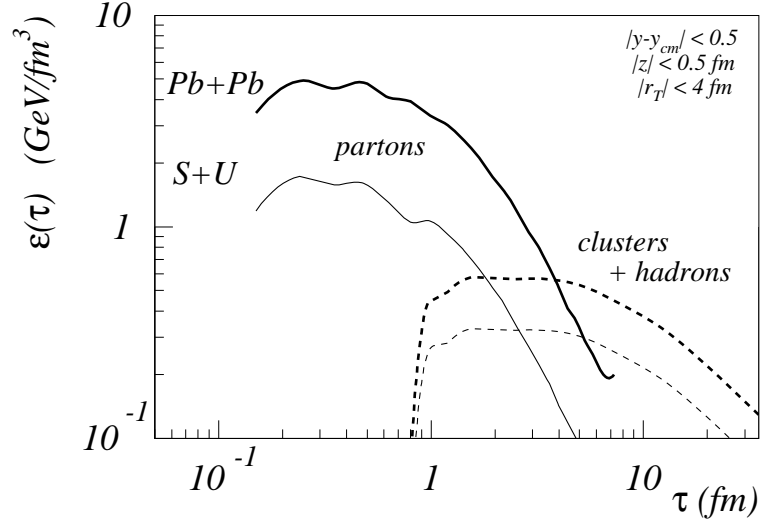

Figure 20: Partonic energy densities for S + U and Pb + Pb in the parton cascade model [41].

On the other hand, the suppression of the ψ' state is so strong already in S + U collisions that its description requires comover suppression. This is almost trivially obvious from the fact that the $\psi'/(J/\psi)$ ratio is constant in $p + A$ collisions, but drops significantly and progressively with decreasing impact parameter (growing dE_T/dy) in S + U [42]. This has been analyzed by many authors [6, 37, 43, 44] who all reach similar conclusions (see Fig. 21 for an example). Values for the absorption of the ψ' by comovers are obtained in the range $\sigma_{co}^{eff}(\psi') = 8$ mb, which is quite reasonable in view of the rather large size of the ψ' and of its low dissociation threshold. In fact, the authors of [44] argue that the ψ' suppression observed in Pb + Pb is not quite as strong as expected on the basis of the S + U data, and that a regeneration mechanism must be introduced to obtain a good fit of the Pb + Pb data.

4.2 Microscopic Models

Microscopic models of heavy ion reactions allow us to study the question whether the models based on Glauber theory adequately describe the space-time dynamics of these reactions. This includes questions such as whether the time dependence of the comover density and the comover composition is modeled correctly in the Glauber approach and whether the
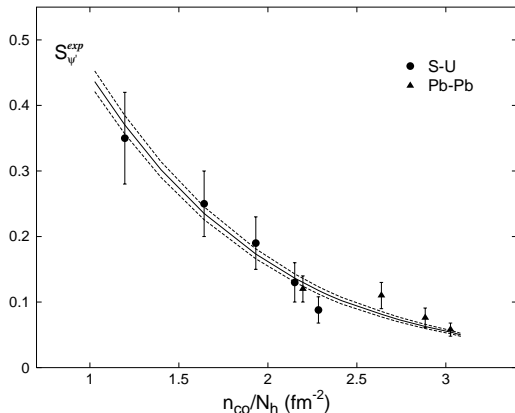
Figure 21: ψ' suppression factor due to absorption on nucleons (with a cross section of 7.3 ± 0.6 mb) as well as absorption on comoving hadrons (with a cross section of about 10 mb), in comparison with the NA38/50 data for the S + U and Pb + Pb systems (from [6]). Comover absorption is required in both systems if one wants to obtain agreement between the data and the Glauber model.

transverse expansion, which is neglected in the Glauber model, plays a significant role.

An extensive study of such questions was recently made by Bratkovskaya and Cassing in the framework of a relativistic hadronic cascade model [45]. The model treats the nuclear collision as a sequence of binary hadronic scatterings with the possibility of a formation time for all newly created particles. When a J/ψ particle is produced in a primary nucleon-nucleon collision, it can interact with another hadron and be dissociated. The authors assume that the $(c\bar{c})$ pair is originally produced in the color octet state and converts to a color singlet J/ψ after $\tau_f = 0.7$ fm/c.

It is useful to begin taking a look at the general space-time structure of a nuclear collision event. The Pb + Pb system differs from higher systems (p + Pb, S + U) in that a significant fraction of J/ψ are formed as color singlets during the time period in which the primary nucleon-nucleon collisions occur. This is so because at the CERN-SPS energy, corresponding to $\gamma_{cm} \approx 9$, the Pb nuclei are Lorentz contracted to a longitudinal width of only about 1.5 fm, exceeding the assumed J/ψ formation length by a factor two. Accordingly, the J/ψ states are on average produced in a more violent environment than in the case for lighter collision systems. The density of comovers is higher in Pb + Pb than in S + U and falls off more slowly.

The Cassing-Bratkovskaya model predicts a peak density of pions which slightly exceeds 1 fm^3 in the Pb + Pb system. A look at our Fig. 15 shows that this pion density is equivalent to a temperature $T \approx 300$ MeV under equilibrium conditions. The same holds true for the peak density of ρ-mesons (about 0.6 fm^{-3}), indicating that the comover abundances may be close to chemical, if not thermal, equilibrium. In other words, the model predicts hadronic comover densities far in excess of those in the range of the presumed validity of the hadronic

phase. The peak densities predicted for the S + U system are lower and correspond to those of a thermal environment of "only" $T \approx 250$ MeV.
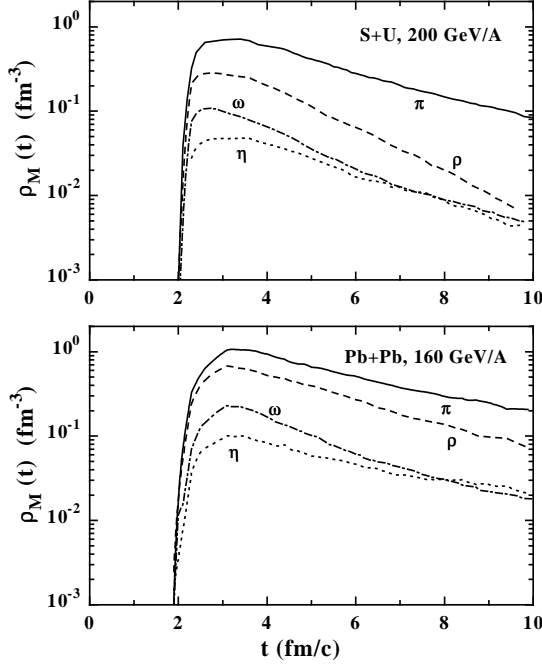

Figure 22: Density of comovers as function of time in the cascade model of Cassing and Bratkovskaya.
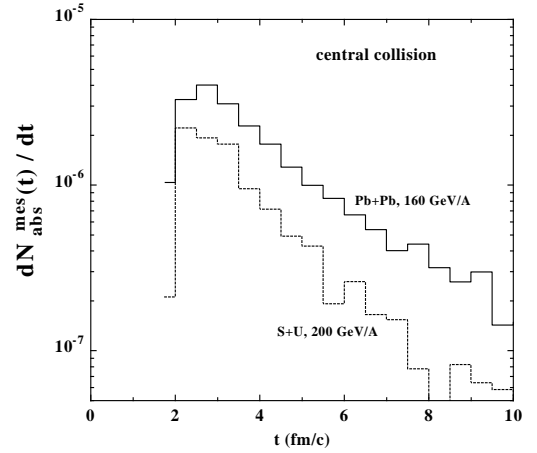
Figure 23: Time distribution of J/ψ absorption events. Most absorptions occur within 3 fm/c after the onset of the reaction.

Figure 23 shows that most of the absorption of J/ψ by comovers in this model occurs within the first 3 fm/c after the initial reaction. At that time ($t = 5$ fm/c in Fig. 23), the comover density still corresponds to a thermal bath at $T \approx 250$ MeV in Pb + Pb and $T \approx 210$ MeV in S + U. This confirms the conclusion reached in Section III.C, that significant dissociation rates for J/ψ in a thermal hadronic environment require temperatures far beyond T_c, at least in the range $T \approx 250 - 300$ MeV.

A second interesting observation is that the time interval during which the J/ψ can be efficiently dissociated is so short that transverse expansion effects can be safely neglected. Hence, the Glauber approximation should be an excellent approximation, and in many respects preferable to microscopic cascade models, because it allows for a more consistent and economical treatment of quantum effects.[7]

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[7]Cascade models can be useful tools for estimating quantities such as the effective optical thickness of the

In conclusion, hadronic cascade models provide useful tools for testing the consistency of the hadronic comover suppression scenario. At present, their results confirm that this scenario has difficulty explaining an observed "anomalous" suppression within the range of validity of the hadronic gas model.

5 Addendum: Developments since June 1997

Since mid-1997, a large number of articles have appeared addressing the implications of the new data from experiment NA50 on J/ψ and ψ' production in Pb + Pb collisions. This activity was triggered by the new results presented by the NA50 collaboration, first at the Quark Matter 1997 conference in Tsukuba [14] and later at the Moriond meeting in March 1998. The new data from the analysis of the 1996 Pb beam run at CERN, reproduced in Figures 6 and 7, revealed a pronounced sudden drop in the ratio of J/ψ production to Drell-Yan pairs in the Pb + Pb system at an impact parameter around 8 fm (corresponding to a nuclear absorption length $L \approx 8$ fm). Furthermore, the new data confirmed the presence of an "anomalous" suppression effect in Pb + Pb, and their much higher statistics leaves hardly any doubt that the effect is real.

After a careful reanalysis of the viability of hadronic comover models Vogt [46] concluded that the data for p + A and S + U collisions leave insufficient room for hadronic comovers to make an explanation of the suppression of J/ψ in Pb + Pb in terms of hadronic absorption viable. This agrees with the results of Kharzeev, et al. [6]. Vogt finds no difficulty with an interpretation of the ψ' data from NA50 in terms of hadronic absorption. Vogt's analysis includes the suppression effects on the feed-down to J/ψ from the excited states χ_c (normally about 30%) and ψ' (12%). The S + U data can be fitted with $\sigma_{c\bar{c}N}^{(\mathrm{abs})} = 7.3$ mb and no comover suppression or with $\sigma_{c\bar{c}N}^{(\mathrm{abs})} = 4.8$ mb and $\sigma_{\mathrm{co}} = 0.67$ mb, 1.6 mb, and 2.5 mb for J/ψ, χ_c, and ψ', respectively. The extrapolation to Pb + Pb then fails to explain the NA50 data, if the comover density is scaled as E_T, as expected on the basis of multiparticle production models, such as the wounded nucleon model. A reasonable agreement with the Pb + Pb data would require a much faster than linear rise of the comover density with E_T, such as $E_T^{5/3}$.[8]

Conclusions similar to those of Vogt are reached by Wong [47], who analyzes in detail

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comoving matter, $\int d\tau v \sigma_{\mathrm{abs}} dN(\tau)/dy$, as input into simplified Glauber calculations. However, one needs to exercise caution since quantum coherence effects that may affect $N(\tau)$ are neglected.

[8]Such a fast rise could be explained by assuming that the comover density relevant for J/ψ absorption scales like other hard QCD processes, as suggested by the arguments presented in [40].

some implications of the competing color singlet and octet production mechanisms for the J/ψ and ψ'. Martins and Blaschke [48], based on their hadronic charm exchange model [27], also rule out the hadronic absorption scenario.

On the other hand, Dias de Deus and Seixas [49] argue that a very simple, schematic model of comover suppression fits the NA50 data in a roundabout way; however, their model does not explain the apparently different behavior of the low- and high-E_T regions in the Pb + Pb system. Frankel and Frati [50] argue that the observed suppression can be explained in terms of the energy loss of partons (gluons) as they travel through the colliding nuclei. Apart from being unable to describe a sudden drop in the J/ψ yield, it is not entirely clear how this model can escape the stringent limits on energy loss set by the Drell-Yan phenomenology.

A number of authors have quantitatively investigated the effectiveness of various absorption mechanisms for J/ψ in dense matter. The reaction $\pi + J/\psi \rightarrow \psi' + \pi$ in the presence of a thermal pion gas was studied in two articles. Chen and Savage [52] analyzed the reaction within the framework of chiral perturbation theory, finding a very small average cross section of order 0.01 mb at thermal energies. Sorge, Shuryak, and Zahed [51] made the assumption that the $\pi\pi$ spectrum in the decay $\psi' \rightarrow J/\psi\pi\pi$ is dominated by a scalar resonance. In their model the ratio $\psi'/(J/\psi)$ can chemically equilibrate to the value 0.05 observed in the highest E_T bins of the Pb + Pb and S + U data, if the mass of the scalar resonance drops at high density, enhancing the reaction rate. Shuryak and Teaney [53] calculated the reaction $\pi + J/\psi \rightarrow \eta_c + \rho$, but found that the thermal rate was only about 10^{-3} $(\text{fm}/c)^{-1}$.

Several authors [46, 48, 54] have emphasized that the major part, if not all, of the additional suppression of the J/ψ yield seen in the Pb + Pb data can be understood as virtually complete absorption of the fraction (about 30%) of J/ψ which originates from decays of χ_c. The deconfinement scenario favors this explanation, because the χ_c disappears almost immediately at T_c, while the J/ψ survives to higher temperatures. This picture is supported by recent, improved lattice calculations [56]. QCD-based calculations of the absorption cross sections of J/ψ and χ_c on nucleons also show the higher vulnerability of the χ_c state at low energies [57]. Gerland et al. [55] argue that the stronger absorption of χ_c is already needed to explain the p + A data and hence must play a significant role in any explanation of the heavy ion data for J/ψ suppression.

The picture of J/ψ suppression by QCD strings has been embraced with considerable enthusiasm lately. Following [21], Geiss et al. [58] simulated this effect within the HSD model discussed in section 4.2. They found reasonable agreement with the NA50 data by setting the string radius to 0.2–0.3 fm and assuming that any J/ψ immediately dissociates

when entering a string. With a string area density of about $3/\text{fm}^2$, this implies that at least half the transverse cross section of the reaction volume in Pb + Pb is filled with QCD strings. A related, but more schematic study by Nardi and Satz [59] concludes that the sudden drop in the J/ψ yield observed by NA50 coincides with the point where QCD strings become so dense that large overlapping clusters develop for an assumed string radius of 0.2 fm. In this picture the drop is associated with a color percolation phase transition which causes the disappearance of the χ_c state. The existence of such a phase transition was also emphasized by Braun, Pajares, and Ranft [60], who studied the percolation of strings analytically as well as by simulations. Much earlier studies of J/ψ absorption reached similar conclusions [20].

6   Summary and Outlook

Have heavy quark bound states, the J/ψ and ψ', fulfilled their promise as probes of the structure of dense hadronic matter formed in relativistic heavy ion collisions? Although the final answer to this question is still extant, several conclusions can be drawn now. Clearly, the original expectation that any significant reduction of J/ψ formation in nuclear collisions would signal the creation of a quark-gluon plasma had to be revised. We now understand that the suppression observed with light nuclear projectiles at the CERN-SPS is a smooth extrapolation of the less-than-linear target mass dependence found in $p + A$ reactions. We also know a credible mechanism for this suppression effect: the color-octet formation model. However, it needs to be stressed that we have almost no direct evidence for the color-octet $(c\bar{c})$ state, and we lack a quantitative description of the transition of the color-octet pair to a color-singlet charmonium state from first principles. For these reasons, the color-octet model must be considered a reasonable, even likely, explanation of the suppression from $p + A$ to S + U reactions, but it is far from being firmly and quantitatively established.

The available experimental evidence, combined with theoretical arguments, points toward the emergence of novel J/ψ suppression mechanism in Pb + Pb collisions. Because no experiments with projectiles between ^{32}S and ^{208}Pb have been done, it is unclear whether this new mechanism sets in suddenly or gradually. However, the magnitude of the effect is such that it can hardly be explained as the absorption of color-singlet J/ψ mesons on hadronic comovers, unless one wants to invoke the existence of hadrons, as we know them, at densities exceeding 1 fm^{-3}. Moreover, a sizable comover suppression effect should be visible already in S + U collisions, if the "anomalous" suppression seen in the Pb + Pb system could be accounted for by absorption on hadrons. On the basis of the existing experimental evidence,

the structure of comoving matter has to undergo a significnt change between S + U and Pb + Pb, if comover absorption is at the origin of the suppression effect in Pb + Pb collisions.

In the following, we enumerate some weaknesses of the present arguments regarding the "anomalous" J/ψ suppression as evidence for the creation of a quark-gluon plasma in Pb + Pb collisions at CERN. We also list some issues warranting further experimental and theoretical study.

1. The argument that J/ψ suppression in S + U collisions is "normal" but that seen in Pb + Pb collisions has a novel, "abnormal" component relies critically on the precision of the extrapolation of the $p+A$ results. As the advocates of hadronic comover models have pointed out, the Pb + Pb data would lose their cogency if the new suppression mechanism were found to set in gradually as the nuclear projectiles get heavier. How large is the error in the value $\sigma_{c\bar{c}N}^{(\text{abs})} = 7.3$ mb, i.e. how uncertain is the extrapolation of the Gerschel-Hüfner line to nuclear systems? The extrapolation from $p + A$ to nuclear collisions depends critically on the assumption that the time for color neutralization is sufficiently long, so that the $(c\bar{c})$ pair remains in the color octet state while the nuclear matter sweeps by. Can one really exclude $\sigma_{c\bar{c}N}^{(\text{abs})} = 5$ mb if all errors in the $p + A$ data and the extrapolation to nuclear systems are taken into account?

2. The viability of the hadronic comover model critically depends on the value of the absorption cross section for J/ψ on hadrons. It is possible to obtain a safe upper bound for $\sigma_{\psi h}^{(\text{abs})}$ in a hadronic gas of imprecisely known composition? A systematic study of J/ψ absorption within effective meson theories including mesons containing heavy quarks would be extremely useful. Another important study would include heavier hadronic states in the operator-product expansion apaproach to J/ψ absorption. Since $\langle h|\alpha_s E^2|h\rangle \sim m_h^2$, heavier mesons can act as ample sources of gluons that may excite and dissociate the J/ψ. (The higher dissociation efficiency of ρ-mesons is also visible in the D-meson exchange model, see section III.C.) For example, it would be interesting to explore the J/ψ dissociation rate in a resonance gas with a Hagedorn-type excitation spectrum.

3. An important issue that has not been investigated extensively is the problem of χ_c absorption [61]. Clearly, $\sigma_{\chi h}^{\text{abs}}$ must lie somewhere between the cross sections for J/ψ and ψ' absorption. The question is if and where "anomalous" χ_c absorption sets in. Can hadronic comover absorption of χ_c in S + U be ruled out? Can the "anomalous"

absorption of J/ψ in Pb + Pb be attributed entirely to simiarly "anomalous" absorption of the component due to feed-down from χ_c? Vogt's extensive study [46] seems to indicate that it is difficult to find a scenario including feed-down from χ_c to J/ψ that is compatible with all existing heavy ion data.

4. There are many improvements that could be made to microscopic models of J/ψ suppression. Including a realistic energy dependence of $\sigma_{\psi h}^{(\text{abs})}$, similar to the one found in section III.C, would probably lead to an effective nonlinear dependence of the comover effect on dN/dy, because higher comover densities are usually correlated with increased average kinetic energy per particle.

   Because the p_T-spectrum of J/ψ-suppression will be used as a tool for discriminating between different mechanisms, a study of final state changes in the p_T-distribution of J/ψ, e.g. due to elastic scattering on comovers, would be interesting. In more general terms, microscopic transport models could be used to systematically assess kinematical effects that can invalidate the Glauber approximation in the context of J/ψ production, in particular, effects of transverse expansion and changes in the p_T distribution of the J/ψ by elastic scattering in the medium [62]. Both effects increase for heavier nuclei.

5. Desirable improvements in the experiment include a reduction of statistical and systematic errors by taking more data. Higher statistics data would also permit the determination of the J/ψ suppression factor in a variety of E_T bins and over a range of values of p_T, the transverse momentum of J/ψ. A systematic study of the mass dependence of the effect for several symmetric systems from S + S to Pb + Pb would be especially important, so that the specific predictions of the deconfinement model can be checked against the data. An excitation function would also be of interest, but due to the steep beam-energy dependence of J/ψ production it may be difficult to eliminate systematic errors with the required accuracy. As pointed out by Kharzeev and Satz [57], the inverse kinematic experiment (Pb + S) would allow for a test of formation time effects on the suppression factor.

6. The comparison of the experimental data with theoretical models will be facilitated, if the data are plotted against measured variables (such as E_T) rather than model dependent derived quantities (such as L). Of course, a detailed understanding of the experimental trigger conditions is indispensible for a meaningful comparison between experiment and theory.

A clear set of ground rules needs to be established for the data analysis as well as for the comparison with theory. If the ratio between J/ψ yield and Drell-Yan (DY) background is plotted, only those mass ranges should be included where the DY process can be unambiguously identified. It makes little sense to include the region under the J/ψ itself, where the DY pairs make only a small contribution.

Theorists also need to take a careful look at the soundness of their models. While microscopic transport models may have the appearance of great realism, their quantitative results can depend on many implicit model assumptions that are difficult to analyze. Thus, more schematic models will continue to provide useful analytic tools. However, blind or ad-hoc parameter fitting can easily confuse the situation. Parameter values must be related to other experimental data, providing support for the fit, and should not exceed the limits of applicability of the schematic model itself. In some cases, it may be possible to check by *ab initio* QCD calculations whether a certain parameter set makes sense. As the picture of J/ψ suppression in nuclear reactions takes on a firmer shape, models must increasingly be judged by their ability to predict new observables quantitatively.

Finally, with the experimental emphasis shifting from the CERN-SPS to RHIC in the years ahead, it may be possible to observe not only the nuclear suppression of the J/ψ, but also the suppression of the Υ in heavy collision systems. Parton cascade models predict initial energy densities of the order of $50\ \mathrm{GeV/fm^3}$, far exceeding the deconfinement threshold of the $(b\bar{b})$ ground state. The advantage of the Υ would be that its suppression in $p+A$ collisions is much weaker than that of the J/ψ and, therefore, its almost complete suppression in a deconfined phase would be quite spectacular.

Acknowledgements:

Appendix A: J/ψ Dissociation in a Medium

Here we consider the dissociation of a J/ψ state in a medium within the Bhanot-Peskin formalism. The fundamental idea is that the dissociation is initiated by the absorption of a gluon from the medium which excites the color-singlet $(c\bar{c})$ pair into a color-octet continuum state $|(c\bar{c})^{(8)}, \epsilon\rangle$, where ϵ denotes the final state energy. The S-matrix element is

$$S_{fi} = \frac{1}{i\hbar} \int_{-\infty}^{\infty} dt \langle (c\bar{c})^{(8)}\epsilon | g\vec{r} \cdot \vec{E}(r,t) | (c\bar{c})^{(1)} \rangle. \tag{26}$$

We will make use of the dipole approximation assuming that the dissociation is dominated by long-range gluonic interactions. In the heavy quark limit, the binding energy of the $(c\bar{c})^{(1)}$ state is $\frac{4}{9}\alpha_s^2 m_Q$, whereas its inverse size is $a^{-1} = \frac{2}{3}\alpha_s m_Q$. When $\alpha_s \ll 1$, the characteristic wavelength of a gluon that is energetically capable of dissociating the J/ψ state is larger than the J/ψ radius. This approximation is somewhat marginal for the J/ψ state; it should be very good for the Υ. After performing the SU(3)-color algebra, the dipole matrix element is given by

$$\frac{g}{\sqrt{6}} \int d^3r \, \varphi_8^{\epsilon}(r) \vec{r} \varphi_1(r), \tag{27}$$

where $\varphi_1(r)$ and $\varphi_8^{\epsilon}(r)$ are the spatial wavefunctions of the singlet and octet states, respectively. We neglect recoil effects, and we use a Coulombic 1s-wavefunction for the singlet state and a plane wave for the octet state. The resulting S-matrix then takes the form

$$S_{fi} = \frac{g}{\sqrt{6}} \sqrt{\frac{\pi a^3}{V}} 32 a^2 \frac{\vec{E}^a(\omega) \cdot \vec{p}}{(1 + p^2 a^2)^3}. \tag{28}$$

Here $\vec{E}^a(\omega)$ is the spectral distribution of gluon fields in the medium at frequency $\omega = \epsilon_8 - \epsilon_1$, a is the Bohr radius of the $(c\bar{c})^{(1)}$ state, $\vec{p}$ is the relative momentum of the $(c\bar{c})^{(8)}$ pair, and V denotes the quantization volume. Integrating over $\vec{p}$ and assuming color neutrality of the medium, i.e.

$$\langle E_i^a(\omega) E_j^b(\omega) \rangle = \frac{1}{24} \delta_{ij} \delta_{ab} \langle |E(\omega)|^2 \rangle, \tag{29}$$

where i, j denote spatial vector and a, b color indices, we finally obtain the transition probability

$$P_{if} = \frac{2}{3} \pi \alpha_s a^2 \langle |E(\omega)|^2 \rangle. \tag{30}$$

Since $\langle |E(\omega)|^2 \rangle$ is proportional to the total interaction time T_{fi}, the result (30) corresponds to a constant dissociation rate.

The color-electric power density $\frac{1}{T_{fi}} \langle |E(\omega)|^2 \rangle$ of the medium can be evaluated analytically for a variety of simple models. One such model is a dilute gas of color charges with a given

density and momentum distribution. Denoting the Casimir of the color charge by Q^2, such a model yields

$$\frac{1}{T_{fi}}\langle |E(\omega)|^2\rangle \approx \frac{\pi}{2}\alpha_s Q^2\tilde{\rho}(\omega),\tag{31}$$

where $\tilde{\rho}(\omega)$ is a ω-dependent, weighted average of the density of charges in the medium.

If the medium is represented as a dilute gas of hadrons, the power spectrum is related to the gluon structure functions of the hadrons $G_h(x)$ where $x = \omega/p_h$, in the spirit of the calculation of Kharzeev and Satz [24]. Given a momentum spectrum of hadrons $f(p_h)$ and chemical abundances of different hadrons, the power spectrum can be evaluated.

Crude estimates of the dissociation rate in these two models yield rather long survival times of the J/ψ in any realistic hadronic medium. For example, counting the valence quarks in mesons and combining (30) with (31), would yield a dissociation rate

$$\Gamma_{\text{dis}} \approx \frac{8}{9}\pi^3\alpha_s^2 a^2\tilde{\rho}_{\text{meson}},\tag{32}$$

where only sufficiently energetic mesons are counted.

Appendix B: Matrix Elements for Charm Exchange

In this Appendix we list the explicit expressions for the invariant matrix elements for the charm exchange reactions described by the Feynman diagrams in Fig. 13. We begin with the two diagrams for J/ψ absorption on pions (diagrams 13a,b). Averaging over initial and summing over final isospin and spin states, we obtain

$$\overline{|M_a|^2} = \frac{8}{3}g_{\psi DD}^2 g_{\pi DD^*}^2 \frac{\left(m_\pi^2 - \frac{(p_\pi \cdot p_{D^*})^2}{m_{D^*}^2}\right)\left(m_D^2 - \frac{(p_\psi \cdot p_D)^2}{m_\psi^2}\right)}{(q^2 - m_D^2)^2}\tag{33}$$

where $q = p_\pi - p_{D^*} = p_D - p_\psi$ is the momentum transfer of the reaction. Introducing the Mandelstam variables s, t and u, and the shorthand notation $t' = t - m_D^2$, we obtain

$$\overline{|M_a|^2} = \frac{1}{6}g_{\psi DD}^2 g_{\pi DD^*}^2 \frac{1}{t'^2}\left(4m_\pi^2 - m_{D^*}^2 + 2t' - \frac{t'^2}{m_{D^*}^2}\right)\left(4m_D^2 - m_\psi^2 + 2t' - \frac{t'^2}{m_\psi^2}\right).\tag{34}$$

The result for the diagram (13b) is identical, i.e.

$$\overline{|M_b|^2} = \overline{|M_a|^2}.\tag{35}$$

The contributions of the diagrams (13c,d) need to be added coherently, because they lead to identical final states. However, their results are related by crossing symmetry ($t \leftrightarrow u$).

After somewhat lengthy algebra one finds:

$$
\overline{|M_c + M_d|^2} = \frac{1}{18}g_{\psi DD}^2 g_{\rho DD}^2 \left[\frac{1}{t'^2}\left(4m_D^2 - \frac{(m_\rho^2 - t')^2}{m_\rho^2}\right)\left(4m_D^2 - \frac{(m_\psi^2 - t')^2}{m_\psi^2}\right) \right.
$$
$$
+\frac{1}{u't'}\left(2s - 4m_D^2 - \frac{(m_\rho^2 - t')(m_\rho^2 - u')}{m_\rho^2}\right)\left(2s - 4m_D^2 - \frac{(m_\psi^2 - t')(m_\psi^2 - u')}{m_\psi^2}\right)
$$
$$
\left. + (t' \leftrightarrow u') \right] \tag{36}
$$

where again $t' = t - m_D^2$, $u' = u - m_D^2$. Of course, the crossing related diagrams (13c,d) yield the same contribution to the total absorption cross section, with an additional interference term.

Finally, for completeness, we give the expression for the decay width of $D^* \to D\pi$:

$$
\Gamma_{D^*} = \frac{g_{\pi DD^*}^2}{8\pi}\frac{p_\pi^3}{m_{D^*}^2}. \tag{37}
$$

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